Evidence for the Collective Nature of the Reentrant Integer Quantum Hall States of the Second Landau Level

N. Deng[1], A. Kumar[1], M.J. Manfra[1,2] , L.N. Pfeiffer[3], K.W. West[3], and G.A. Csáthy[1]
[1] *Department of Physics, Purdue University, West Lafayette, IN 47907, USA*
[2] *Birck Nanotechnology Center, School of Materials Engineering and School of Electrical and Computer Engineering, Purdue University, West Lafayette, IN 47907, USA*
[3] *Department of Electrical Engineering, Princeton University, Princeton, NJ 08544*

(Dated: October 20, 2011)

We report an unexpected sharp peak in the temperature dependence of the magnetoresistance of the reentrant integer quantum Hall states in the second Landau level. This peak defines the onset temperature of these states. We find that in different spin branches the onset temperatures of the reentrant states scale with the Coulomb energy. This scaling provides direct evidence that Coulomb interactions play an important role in the formation of these reentrant states evincing their collective nature.

The second Landau level (SLL) of the two-dimensional electron gas (2DEG) is astonishingly rich in novel ground states [1–3]. Recent experiments [3–9] suggest that there are both conventional [10, 11] as well as exotic fractional quantum Hall states (FQHSs) [12, 13] in this region. The study of the latter has enriched quantum many-body physics with numerous novel concepts such as paired composite fermion states with Pfaffian correlations, non-Abelian quasiparticles [12–19], topologically protected quantum computing [20], and established connections between the 2DEG and p-wave superconductivity in Sr_2RuO_4 and fermionic atomic condensates.

The eight reentrant integer quantum Hall states (RIQHSs) form another set of prominent ground states in the SLL [1]. The transport signatures of the RIQHSs are consistent with electron localization in the topmost energy level [1]. However, the nature of the localization is not yet well understood. Depending on the relative importance of the electron-electron interactions, the ground state can be either an Anderson insulator or a collectively pinned electron solid.

FQHSs owe their existence to the presence of the interelectronic Coulomb interactions [10, 11]. Since FQHSs and RIQHSs alternate in the SLL, it was argued that Coulomb interactions must be important and, therefore, the RIQHSs in the SLL must be electron solids [1]. Subsequent density matrix renormalization group [21] and Hartree-Fock calculations [22] also favored the electron solid picture and predicted the solid phase similar to the Wigner crystal, but having one or more electrons in the nodes of the crystal [22]. Recently reported weak microwave resonances in one such RIQHS are suggestive of but are far from being conclusive on the formation of a collective insulator [23]. Our understanding of the RIQHSs in the SLL, therefore, is still in its infancy and the collective nature of these states has not yet been firmly established.

We report a feature in the temperature dependent magnetoresistance unique to the the RIQHSs in the SLL, a feature which is used to define the onset temperature of these states. The scaling of onset temperatures with the Coulomb energy reveals that Coulomb interactions play a central role in the formation of RIQHSs and, therefore, these reentrant states are exotic electronic solids rather than Anderson insulators. We also report an unexpected trend of the onset temperatures within each spin branch. This trend is inconsistent with current theories and can be understood as a result of a broken electron-hole symmetry. Explaining such a broken symmetry of the RIQHSs is expected to impact our understanding of a similar asymmetry of the exotic FQHSs of the SLL, including the one at $\nu = 5/2$.

We performed magnetotransport measurements on a high quality GaAs/AlGaAs sample of density n=3.0×10^{11}cm^{-2} and of mobility μ=3.2×10^7cm^2/Vs. Earlier we reported the observation of a new FQHS at $\nu = 2 + 6/13$ in this sample [3]. The sample is immersed into a He-3 cell equipped with a quartz tuning fork viscometer used for B-field independent thermometry [24].

In the top panel of Fig.1 we show the Hall resistance R_{xy} in the SLL at 6.9 mK. The data reveals numerous FQHSs and it is dominated by the eight RIQHSs. Starting with the states at the highest B-field we label the RIQHSs with $R2a$, $R2b$, $R2c$, and $R2d$ in the lower spin branch of the SLL (i.e. $2 < \nu < 3$) and with $R3a$, $R3b$, $R3c$, and $R3d$ in the upper spin branch (i.e. $3 < \nu < 4$). Here $\nu = nh/eB$ is the Landau level filling factor. RIQHSs have historically been predicted [25] and observed [26, 27] in high Landau levels (i.e. $\nu > 4$). In contrast to the SLL, in high Landau levels there are only four RIQHSs in each Landau level. These states develop at the lowest temperatures around non-integer filling fac-

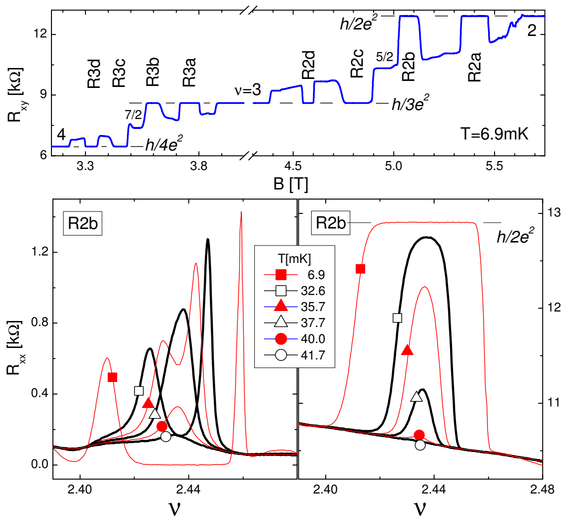

FIG. 1. The Hall resistance of the eight RIQHSs in the SLL at 6.9 mK (top panel) and the temperature evolution of the RIQHS labeled *R2b* (bottom panels). Numbers mark Landau level filling factors of importance.

tors and yet their Hall resistance R_{xy} is quantized to a nearby integer plateau [1].

Because of the delicate nature of the RIQHS in the SLL [1–6, 28–33] there is only scarce information available on their temperature dependence [28, 29, 32]. The lower panels of Fig.1 show the details of the evolution of the longitudinal resistance R_{xx} and R_{xy} of *R2b* with temperature T. The $R_{xx}(B)|_{T=6.9\mathrm{mK}}$ curve has a wide zero flanked by two sharp spikes. As the temperature is raised, the spikes in R_{xx} persist but they move closer to each other and the width of the zero decreases. At 32.6 mK $R_{xx}(B)$ does still exhibit the two spikes but instead of a zero it has a non-zero local minimum. The location in B-field of this minimum is T-independent and it defines the center $\nu_c = 2.438$ of the *R2b* state. At 35.7 mK the two spikes of $R_{xx}(B)$ have moved closer to each other and between them there is still a local minimum, albeit with a large resistance. A small increase in T of only 2 mK leads to a qualitative change. Indeed, in contrast to curves at lower T, $R_{xx}(B)|_{T=37.7\mathrm{mK}}$ exhibits a single peak only. As the temperature is further raised, this single peak rapidly decreases until it merges into a low resistance background. Simultaneously with the described changes of R_{xx}, R_{xy} evolves from the quantized value $h/2e^2$ to its classical value $B/ne = h/\nu_c e^2$.

The behavior seen in Fig.1 can be better understood by measuring T-dependence at a fixed ν. In Fig.2 we show R_{xy} versus T near the center ν_c of the various RQIHSs. It is found that R_{xy} assumes the classical Hall resistance at high temperatures and it is quantized to $h/2e^2$ or $h/3e^2$ at the lowest temperatures. Since 80% of the change in R_{xy} between these two values occurs over only 5 mK, this change is very abrupt and it clearly separates the RIQHS at low T from the classical gas at high T. We interpret the inflextion point in R_{xy} versus T as being the onset temperature T_c of the RIQHS. For reliable measurements in the vicinity of T_c the temperature is swept slower than 10 mK/hour.

A transition from the classical Hall value to a quantized R_{xy} with decreasing T is observed not only for the RIQHSs in the SLL but also in the vicinity of any developed integer or fractional quantum Hall state and it is due to localization in the presence of a B-field. As seen in Fig.2, the $R_{xx}(T)|_{\nu=\mathrm{fixed}}$ curves for the RIQHSs are non-zero at high T, they vanish at low T, and they exhibit a sharp peak at the onset temperature T_c defined above. In contrast, $R_{xx}(T)|_{\nu=\mathrm{fixed}}$ of a quantum Hall state changes monotonically, without the presence of a peak. The sharp peak in $R_{xx}(T)|_{\nu=\mathrm{fixed}}$ is, therefore, a signature of localization *unique* to the RIQHSs in the SLL and the peak temperature can be used as an alternative definition for the onset temperature T_c.

Fig.3 represents the stability diagram of the RIQHSs in the ν^*-T plane. As described earlier, at a given ν the RIQHSs develop below the peak present in the $R_{xx}(T)|_{\nu^*=\mathrm{fixed}}$ curve. Such peaks are shown in Fig.2 for $\nu^* \approx \nu_c^*$, but similar peaks are also present for nearby filling factors (not shown). Open symbols in Fig.3 are the peak temperatures T_c as plotted against ν^*. Similarly, the RIQHSs develop between the spikes of the $R_{xx}(\nu)|_{T=\mathrm{fixed}}$ curves. Such spikes are shown for *R2b* in the lower panels of Fig.1. The filling factors ν^* of the spikes for each RIQHS measured at a given temperature

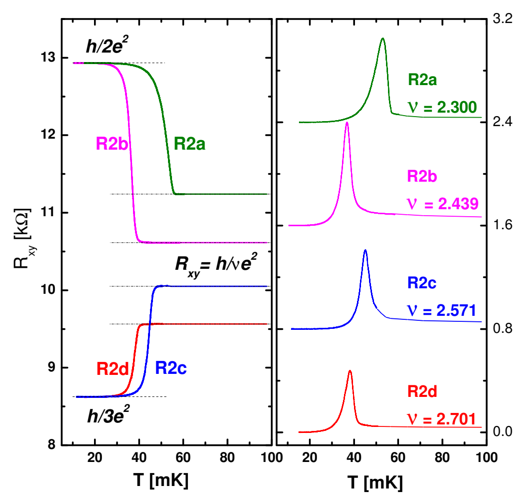

FIG. 2. The evolution of the magnetoresistance of RIQHSs of the lower spin branch with temperature near the center ν_c of each RIQHS. For clarity, $R_{xx}(T)$ curves have been shifted vertically by 0.8 kΩ.

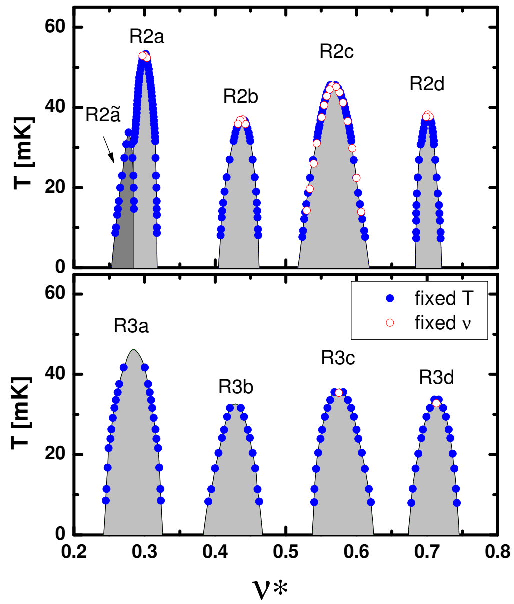

FIG. 3. The phase boundaries of the eight RIQHS in the SLL in the ν^*-T plane. The RIQHSs are stable within the shaded areas. Below 33 mK the $R2a$ state has a split-off state labeled $R2\tilde{a}$.

are marked with closed symbols in Fig.3. The excellent overlap of the two data sets in Fig.3 shows that the two definitions used above self-consistently define the stability boundary of each RIQHS. The shaded areas within each boundary of Fig.3 represent the RIQHSs. FQHSs can develop only outside these shaded areas. The locations $\nu^*_{\rm high}$ and $\nu^*_{\rm low}$ of the spikes of the $R_{xx}(\nu)|_{T=\rm fixed}$ curve measured at the lowest $T = 6.9$ mK of our experiment are listed in Table I.

We note that the $R2a$ state is different from the rest of the RIQHSs as it splits into two RIQHSs with a decreasing temperature. Such a split is signaled by an R_{xy} deviating from $h/2e^2$ as well as a non-zero R_{xx} in the vicinity of $\nu = 2 + 2/7$ and it has already been reported in Ref.[2]. The split-off RIQHS is marked as $R2\tilde{a}$ and with a darker shade in Fig.3. We note that our data is similar to that in Refs.[1] in that the Ria, $i = 2, 3$ is the most stable state. Other studies find the $R2c$ state to be the most stable of RIQHSs [4–6, 23, 28–33].

Each stability boundary shown in Fig.3 can be fitted close to their maxima with a parabolic form $T_c(\nu^*) = T_c(\nu^*_c) - \beta(\nu^* - \nu^*_c)^2$. The obtained parameters are listed in Table I. T_c obtained from the fit is within 1 mK from the peak temperature obtained from Fig.2. The centers ν^*_c of the RIQHSs in the upper spin branch are in excellent agreement with the earlier reported values [1]. Those of the upper spin branch, however, have not yet been documented and they differ significantly from those of the lower spin branch. Indeed, $\nu^*_{c,R2\alpha} \neq \nu^*_{c,R3\alpha}$ for $\alpha = a, b, c$, or d, the difference being the largest for the states a and d. Such a difference is not expected from the theory [21, 22] and we think it is due to the interaction of the electrons in the topmost Landau level with those in the filled lower levels. Furthermore, we establish that the centers ν^*_c of RIQHSs in both spin branches obey particle-hole symmetry, as assumed by the theory [21, 22]. In short $\nu^*_{c,Ria} = 1 - \nu^*_{c,Rid}$ and $\nu^*_{c,Rib} = 1 - \nu^*_{c,Ric}$ for $i = 2, 3$, relations which hold within our measurement error for the filling factor of ± 0.003.

In contrast to the centers of the RIQHSs, other parameters of the RIQHSs from Table I. do not obey particle-hole symmetry. These parameters are the maximum onset temperatures $T_c(\nu^*_c)$, the fit parameter β describing the curvature of the stability diagrams near $T_c(\nu^*_c)$, and the widths $\Delta\nu = \nu^*_{\rm high} - \nu^*_{\rm low}$ of the stability regions of the RIQHSs at $T = 6.9$ mK. Indeed, particle-hole symmetry within a spin branch would imply a scaling of T_c with the Coulomb energy E_C and, therefore, with $1/\sqrt{\nu}$. Here $E_C = e^2/\epsilon l_B$ and $l_B = \sqrt{\hbar/eB}$ is the magnetic length. From such a scaling one expects a decreasing $T_c(\nu^*_c)$ with ν^*_c. The increasing trend of T_c with ν^* across $\nu^* = 1/2$ shown in Fig.4a clearly does not obey such a scaling [29] and, therefore particle-hole symmetry assumed in current theories [21, 22] is violated. The non-monotonic dependence of T_c on ν^*_c shown in Fig.4a is therefore at odds with the sequence of the one- and two-electron bubbles suggested [21, 22] and could be a consequence of either Landau level mixing, disorder, or finite thickness effects. Understanding the origin of this broken symmetry is most likely related to and, therefore, is expected to impact our understanding of the similar symmetry breaking of the Pfaffian and anti-Pfaffian construction for the $\nu = 5/2$ FQHS [34–41].

We find that the onset temperatures $T_c(\nu^*_c)$ in the higher spin branch are consistently smaller than those in the lower spin branch. We notice, however, a startlingly similar non-monotonic dependence within each spin branch. A particularly revealing plot is that of the reduced onset temperatures $T_c(\nu^*_c)/E_C$ against the filling factor ν^*_c. As shown in Fig.4b, there is a surprizingly good collapse of $T_c(\nu^*_c)/E_C$ for the different spin branches. This collapse shows that Coulomb interactions

TABLE I. Parameters extracted from the ν^*-T diagram. T_c and β are in units of mK.

	$R2a$	$R2b$	$R2c$	$R2d$	$R3a$	$R3b$	$R3c$	$R3d$
ν^*_c	0.300	0.438	0.568	0.701	0.284	0.429	0.576	0.712
$T_c(\nu^*_c)$	53.0	37.1	45.8	38.0	46.3	32.3	36.1	33.8
$\beta \times 10^{-4}$	10	3.9	2.4	8.5	2.1	2.0	1.6	2.3
$\nu^*_{\rm high}$	0.317	0.461	0.613	0.719	0.324	0.463	0.621	0.742
$\nu^*_{\rm low}$	0.258	0.407	0.523	0.684	0.245	0.388	0.540	0.677

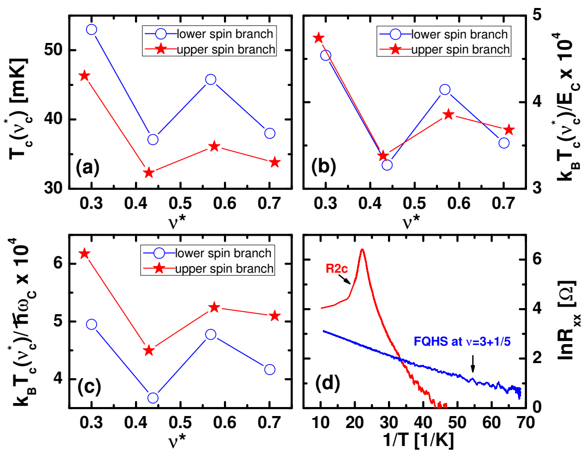

FIG. 4. The variation with the filling factor ν_c^* of the onset temperatures $T_c(\nu_c^*)$ at the center in units of mK (panel a), Coulomb energy (panel b), and cyclotron energy (panel c) of the eight observed RIQHSs in the SLL. Panel d is an Arrhenius plot for $R2c$ and for the $\nu = 3 + 1/5$ FQHS. Lines are guides to the eye.

play a central role in the formation of the RIQHSs in the SLL and provides a first direct evidence that these states reflect collective behavior of the electrons rather than single particle localization. The lack of collapse of $T_c(\nu_c^*)/\hbar\omega_C$ shown in Fig.4c means that $T_c(\nu_c^*)$ does not scale with the cyclotron energy $\hbar\omega_C$.

In a recent study an activated dependence of $R_{xx}(T)$ is found for the $R2c$ state [32]. In our sample we find a significant deviation from such a dependence and, as a consequence, the definition of an activation energy is no longer possible. Fig.4d shows such a plot, together with the activated resistance of a FQHS measured in order to rule out thermometry artifacts. Our data suggest that non-activated behavior might be an inherent property of the RIQHSs.

In summary we find that the scaling of the onset temperatures in different spin branches with the Coulomb energy provides a direct experimental evidence for the collective nature of the RIQHSs of the SLL. The stability diagram we report in the ν^*-T plane reveals several quantitative disagreements with the existing theories such as the lack of paricle-hole symmetry of the onset temperatures within one spin branch.

The work at Purdue was supported by the DOE grant de-sc0006671. L.N.P. and K.W.W. were supported by the Princeton NSF-MRSEC and the Moore Foundation. We benefited from discussions with G. Giuliani and S. Das Sarma.

[1] J.P. Eisenstein *et al.*, Phys. Rev. Lett. **88**, 076801 (2002).
[2] J.S. Xia, *et al.* Phys. Rev. Lett. **93**, 176809 (2004).
[3] A. Kumar *et al.*, Phys. Rev. Lett. **105**, 246808 (2010).
[4] I.P. Radu *et al.*, Science **320**, 899 (2008).
[5] M. Dolev *et al.*, Phys. Rev. Lett. **107**, 036805 (2011).
[6] N. Samkharadze *et al.*, Phys. Rev. B. **84**, R121305 (2011).
[7] M. Dolev *et al.*, Nature **452**, 829 (2008).
[8] R.L. Willett, L.N. Pfeiffer, and K.W. West, Phys. Rev. B **82**, 205301 (2010).
[9] V. Venkatachalam *et al.*, Nature **469**, 185 (2011).
[10] R.B. Laughlin, Phys. Rev. Lett. **50**, 1395 (1983).
[11] J.K. Jain, Phys. Rev. B. **40**, 8079 (1989).
[12] G. Moore and N. Read, Nucl. Phys. B **360**, 362 (1991).
[13] M. Greiter, X.G. Wen, and F. Wilczek, Phys. Rev. Lett. **66**, 3205 (1991).
[14] K. Park, V. Melik-Alaverdian, N.E. Bonesteel, and J.K. Jain, Phys. Rev. B. **58**, R10167 (1998).
[15] N. Read, and E. Rezayi, Phys. Rev. B. **59**, 8084 (1999).
[16] V. Scarola, K. Park, and J.K. Jain, Nature (London) **406**, 863 (2000).
[17] E.H. Rezayi and F.D.M. Haldane, Phys. Rev. Lett. **84**, 4685 (2000).
[18] N. Read and D. Green, Phys. Rev. B. **61**, 10267 (2000).
[19] H. Lu, S. Das Sarma, and K. Park, Phys. Rev. B. **82**, R201303 (2010).
[20] S. Das Sarma, M. Freedman, and C. Nayak, Phys. Rev. Lett. **94**, 166802 (2005).
[21] N. Shibata and D. Yoshioka, J. Phys. Soc. Jpn. **72**, 664 (2003).
[22] M.O. Goerbig, P. Lederer, and C.M. Smith, Phys. Rev. B. **68**, 241302 (2003); Phys. Rev. B. **69**, 115327 (2004).
[23] R.M. Lewis *et al.*, Phys. Rev. B. **71**, R081301 (2005).
[24] N. Samkharadze *et al.* Rev. Sci. Instr. **82**, 053902 (2011).
[25] A.A. Koulakov, M.M. Fogler, and B.I. Shklovskii, Phys. Rev. Lett. **76**, 499 (1996).
[26] M.P. Lilly *et al.*, Phys. Rev. Lett. **82**, 394 (1999).
[27] R.R. Du *et al.*, Solid State Commun. **109**, 389 (1999).
[28] G.A. Csáthy *et al.*, Phys. Rev. Lett. **94**, 146801 (2005).
[29] W. Pan, *et al.* Phys. Rev. B. **77**, 075307 (2008).
[30] C.R. Dean *et al.*, Phys. Rev. Lett. **100**, 146803 (2008).
[31] V. Umansky *et al.* J. Crystal Growth **311**, 1658 (2009).
[32] J. Nuebler *et al.*, Phys. Rev. B. **81**, 035316 (2010).
[33] J. Xia *et al.*, Phys. Rev. Lett. **105**, 176807 (2010).
[34] M. Levin, B.I. Halperin, and B. Rosenow, Phys. Rev. Lett. **99**, 236806 (2007).
[35] S.S. Lee *et al.*, Phys. Rev. Lett. **99**, 236807 (2007).
[36] X. Wan *et al.*, Phys. Rev. B. **77**, 165316 (2008).
[37] H. Wang, D.N. Sheng, and F.D.M. Haldane, Phys. Rev. B. **80**, 241311 (2009).
[38] W. Bishara and C. Nayak, Phys. Rev. B. **80**, 121302 (2009).
[39] M.R. Peterson, K. Park, and S. Das Sarma, Phys. Rev. Lett. **101**, 156803 (2008).
[40] A. Wójs, C. Toke, and J.K. Jain, Phys. Rev. Lett. **105**, 096802 (2010).
[41] E.H. Rezayi and S.H. Simon, Phys. Rev. Lett. **106**, 116801 (2011).